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Tunneling Effects in Tilted Magnetic Fields in *n*-InGaAs/GaAs Structures with Strongly Coupled Double Quantum Wells

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Abstract—The effects of tunneling between two parallel two-dimensional electron gases in *n*-InGaAs/GaAs nanostructures with strongly coupled double quantum wells with a change in the in-plane component of a tilted magnetic field (up to $B_{||} = 9.0$ T) in the temperature range $T = 1.8–70.0$ K are investigated. A nonmonotonic temperature dependence of the inverse quantum lifetime $\tau_q^{-1}(T)$ is obtained from analysis of the dependence of the longitudinal resistance on the parallel component of the tilted magnetic field at fixed temperatures, $\rho_{xx}(B_{||}, T)$. The quadratic portion of this dependence is found to be due to the contribution of inelastic electron–electron scattering. The decrease in the inverse quantum lifetime $\tau_q^{-1}(T)$ at $T > 0.1T_F$ cannot be described within known theories; it seems, it is not related to the processes of electron momentum relaxation.

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1. INTRODUCTION

Tunneling is one of the most striking quantum effects in condensed-matter physics as a whole and in the physics of low-dimensional systems in particular. A number of recent publications [1–8] were devoted to resonant tunneling in two-layer electronic systems containing a two-dimensional electron gas, for example, in structures with two tunnel-coupled double quantum wells (DQWs) with asymmetric scattering. Both transverse [4, 5] and longitudinal [2, 3, 8] transport with a change in the voltage across the gate have been studied. In the first and second cases the dependences of, respectively, the tunneling current (transverse conductance or tunneling differential conductance) and the longitudinal resistance on the gate voltage are measured.

In the transverse-transport technique a change in the voltage across the gate leads to a change in the electron concentrations in the quantum wells. When the Fermi energies in the wells coincide, resonant tunneling occurs and electrons are equiprobably distributed between the two wells. A maximum arises in the dependences of the tunneling current (conductance) on the gate voltage. Tunneling is lacking if the electron concentrations in the wells are significantly different.

In the case of longitudinal transport, when the electron concentrations in the wells significantly differ and tunneling is lacking, we deal with two independent quantum wells connected parallel. If the electron mobilities in the wells are significantly different, the effective resistance of this structure is determined by the well with higher electron mobility (shunting effect). The resonant tunneling of electrons occurs also when the Fermi energies in the wells coincide. In this case, electrons are delocalized between the two wells and their wavefunctions are strongly mixed to form symmetric (S) and antisymmetric (AS) states, separated by the energy gap Δ_{SAS}. The effective resistance of this structure is determined by the total resistance of the two wells. A maximum arises in the dependences of the longitudinal resistance on the gate voltage, as well as in the dependences of the conductance. This phenomenon is referred to as resonance resistance.

Resonance resistance was observed in a number of studies [1–5]; it is explained in terms of the coherent coincidence of tunnel-coupled electron states. The amplitude of the maximum resonance resistance was found to be determined by the ratio of electron mobilities in the wells, i.e., by the transport momentum relaxation times, whereas the shape of the maxima of both the conductance and resonance resistance (peak width) is determined by the broadening of electron states, i.e., the quantum momentum-relaxation times

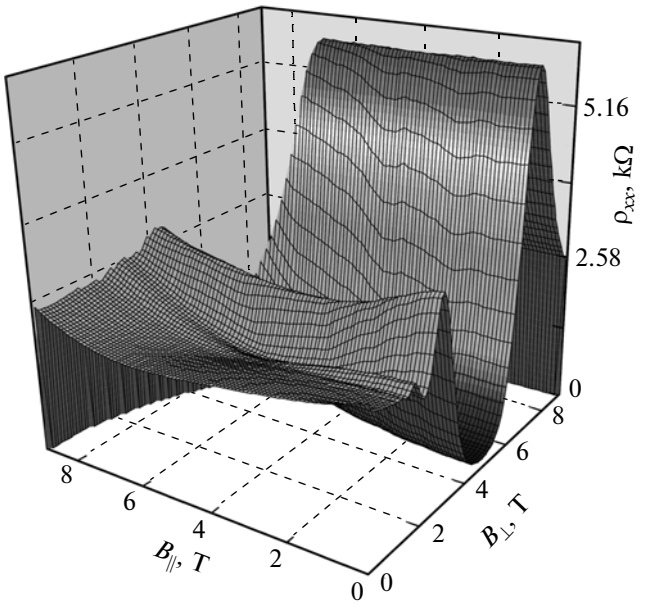

Fig. 1. Dependences of the sample magnetoresistance $\rho_{xx}(B_\perp, B_\parallel)$ (volume pattern) in parallel and perpendicular magnetic fields at the temperature $T = 1.8$ K.

(lifetimes) of electrons. It is well known that these times may significantly differ for selectively doped two-dimensional structures, which exhibit scattering according to the long-range potential of remote impurities. Using these techniques in a certain temperature range, one can obtain the temperature dependences of these relaxation times.

In this study another implementation of the resonance resistance phenomenon is employed. Specifically, resonance resistance is investigated as a function of the parallel component of a tilted magnetic field ($B_\parallel$). The basic concept is as follows. It is assumed, in zero magnetic field, the electron system is in equilibrium, the electron concentrations in the quantum wells are equal, resonance tunneling occurs, and resonance resistance arises. In this case, the Fermi surfaces of the two quantum wells are circles with the same center at the point $k_x = k_y = 0$. A magnetic field directed along k_y leads to the mutual displacement of the centers of Fermi surfaces along k_s by the value $k_B = eB_\parallel b/\hbar$, where b is the distance between the centers of the two quantum wells. As a result, the energy- and momentum-conservation laws are satisfied only in the vicinity of the two intersection points of the Fermi surfaces. The magnetic field oriented parallel to the structure layers (in-plane field), removes the tunnel-coupled states from tunnel resonance. As a result, the resistance in resonance (peak height) decreases with increasing magnetic field. This phenomenon, known as the suppression of resonance resistance by a magnetic field, was observed in recent experiments [2, 3].

The physical reasons for the suppression of resonance resistance and its anisotropy with respect to the angle between the magnetic-field direction and the current through the sample are qualitatively understood at present. It was found that the character of the dependence (line shape) of the resonance resistance on the parallel magnetic field is also determined by the diffusion of electron states, i.e., the quantum lifetime of electrons. Having studied the dependence of the resonance resistance at different temperatures, one can measure the temperature dependences of the quantum lifetime and thus investigate the scattering mechanisms giving rise to broadened energy levels, in particular inelastic electron–electron scattering. It should also be noted that quantitative description of the resonance resistance in a parallel magnetic field has been poorly developed for a number of reasons.

2. EXPERIMENTAL RESULTS AND DISCUSSION

We investigated *n*-InGaAs/GaAs samples with strongly coupled DQWs, in which equilibrium conductance was implemented over subbands of the S and AS states, separated by a tunneling gap. The *n*-InGaAs quantum wells with a width of $d_w = 5$ nm were separated by a GaAs barrier layer with a thickness of $d_B = 10$ nm. The structure was symmetrically doped with Si ($N_D = 10^{18}$ cm^{-3}) in the barriers. The electron-transport and doped regions were located at a distance of (spacer-layer width) $d_x = 19$ nm [7]. The $\rho_{xx}(B, T)$ and $\rho_{xy}(B, T)$ values were measured at $1.8 < T < 70$ K in a tilted magnetic field with detailed scanning of the $(B_\parallel, B_\perp)$ plane within the entire circle from the maximum magnetic field attainable in the solenoid (9.0 T) to zero. To this end, we employed a precision electronically controlled programmable rotator (by Quantum Design firm), which makes it possible to change the angle of magnetic-field rotation with respect to the normal to the layer plane with a step of 0.1° [9]. Then, based on the series of obtained curves $\rho_{xx}(B, T)$ and $\rho_{xy}(B, T)$, we plotted the 3D surfaces $\rho_{xx}(B_\perp, B_\parallel)$ and $\rho_{xy}(B_\perp, B_\parallel)$ using a special interpolation program; the surface obtained at $T = 1.8$ K is shown in Fig. 1. Figures 2 and 3 present the experimental dependences $\rho_{xx}(B_\perp, T)$ and $\rho_{xx}(B_\parallel, T)$. It can be seen that negative magnetoresistance (NMR) occurs in both perpendicular and parallel magnetic fields. It is well known that NMR in $B_\perp$ is related to the suppression of weak-localization quantum effects.

Having analyzed the Shubnikov–de-Haas oscillations in a perpendicular magnetic field at low temperatures ($T \le 8$ K), the Hall effect in weak (R_{H0} ($B_\perp$, T)) and strong ($R_H(B_\perp, T)$) magnetic fields, and the positive magnetoresistance in a perpendicular magnetic field at $T \ge 20$ K, we determined the concentration, mobility, and transport momentum relaxation times of electrons in subbands of the S and AS states and their

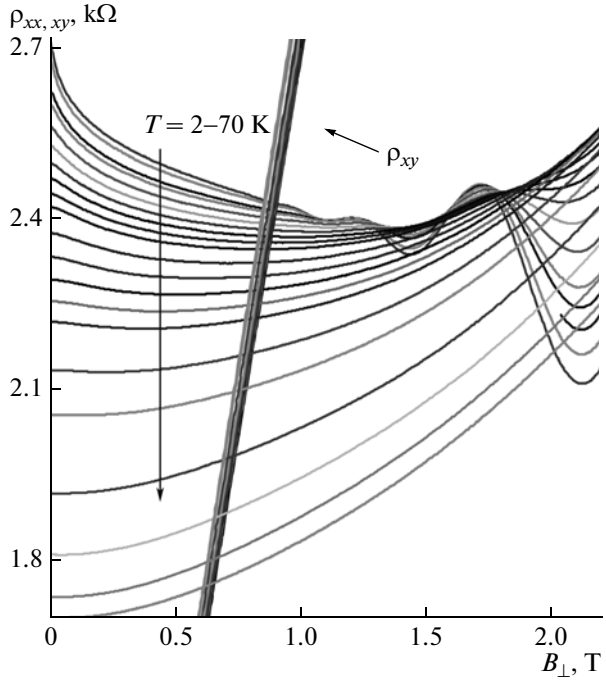

Fig. 2. Dependences $\rho_{xx}(B_\perp, T)$ and $\rho_{xy}(B_\perp, T)$ in a perpendicular magnetic field at temperatures of $T = 2.0$–70.0 K.

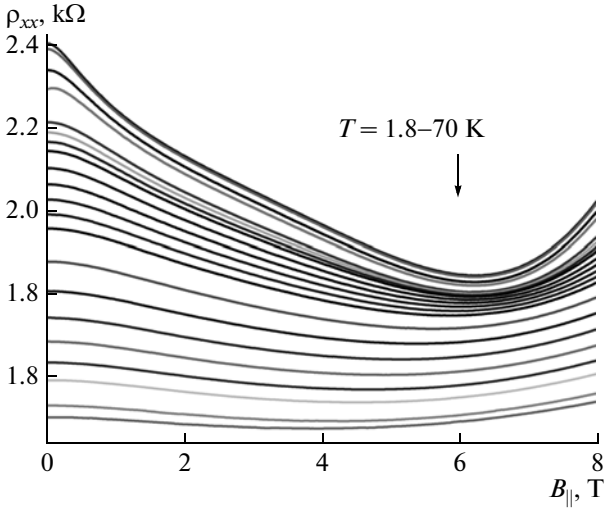

Fig. 3. Dependence $\rho_{xx}(B_\parallel, T)$ for a sample in a parallel magnetic field $B_\parallel$ at $T = 1.8$–70.0 K.

temperature dependences ($n_{1,2}(T)$, $\tau^{tr}_{1,2}(T)$) [7]. The total electron concentration ($n_T = n_1 + n_2 = 2.27 \times 10^{11}$ cm^{-2}) was also determined by analyzing the positions of the peaks and the plateau for the quantum Hall effect. An analysis of the activation dependence of the longitudinal resistance $\rho_{xx}(B_\perp, T)$ in minima in the quantum Hall effect yielded the tunneling-gap value: $\Delta_{SAS} \approx 3.0$ meV [7].

The following expression for the resistance in $B_\parallel \parallel k_y$ was obtained in [2]:

$$\rho_{xx}^{-1}(B_\parallel) - \rho_{off}^{-1} = [\rho(0)^{-1} - \rho_{off}^{-1}] f(B_\parallel / B_c), \qquad (1)$$

where

$$f(x) = 2[(1 + x^2)^{-0.5} - 1]x^{-2}$$

and ρ_{off} is the minimum resistance obtained when DQW is out of resonance (Fig. 3).

The characteristic magnetic field $B_c \ll B_F = 2\pi\hbar/e\lambda_F b$ (λ_F is the Fermi electron wavelength) has the form

$$B_c = \frac{\hbar}{e} \frac{1}{v_F \tau_q b} \sqrt{1 + \left(\frac{\Delta_{SAS}}{\hbar}\right)^2 \left(\frac{\tau_1^{tr} + \tau_2^{tr}}{2} \tau_q\right)}, \qquad (2)$$

where $2\tau_q^{-1} = \tau_{q1}^{-1} + \tau_{q2}^{-1}$ and v_F is the Fermi velocity.

It follows from expression (1) that B_c is the only fitting parameter. Having fitted NMR in weak parallel magnetic fields for temperatures of $T = 1.8$–70.0 K, we obtained the dependence $B_c(T)$ and then worked with expression (2). Note that all electron parameters entering (2) (except for τ_q) were determined from an analysis of the magnetotransport in a perpendicular magnetic field (see above). This approach allowed us to obtain the dependence $\hbar/\tau_q(T)$, which is presented by asterisks in Fig. 4.

The quantum lifetime τ_q is known [4] to be determined by three mechanisms of electron scattering: by ionized impurities ($\tau_q^{e\text{-}imp}$) (independent of T), phonons ($\tau_q^{e\text{-}ph}$), and electrons ($\tau_q^{e\text{-}e}$). The latter two mechanisms are temperature-dependent. Thus, the expression for the temperature dependence of the quantum scattering rate has the form

$$(\tau_q(T))^{-1} = (\tau_q^{e\text{-}imp})^{-1} + (\tau_q^{e\text{-}ph}(T))^{-1} + (\tau_q^{e\text{-}e}(T))^{-1}.$$

In the case of scattering by remote impurities (small-angle scattering in nanostructures with a spacer layer),

$$\frac{\hbar}{\tau_q^{e\text{-}imp}} = \frac{\hbar^2 N_D}{2m^* d_s} \sqrt{\frac{\pi}{2}} n^{-0.5}. \qquad (3)$$

Equation (3) was derived at $4k_F d_s \gg 1$ ($4k_F d_s \approx 10$ for our samples). It was shown in [10] that in the case of electron scattering by phonons the contribution $(\tau_q^{e\text{-}ph})^{-1}$ is comparable with $(\tau_{tr})^{-1}$; therefore, it can be found from the dependences of the mobility on T. The quantum lifetime in selectively doped structures (with a spacer layer) is much shorter than the transport relaxation time; therefore, the scattering rates are described by the relation $(\tau_q^{e\text{-}ph})^{-1} \approx (\tau_{tr})^{-1} \ll (\tau_q^{e\text{-}e})^{-1}$. In addition, scattering by phonons is known to be pro-

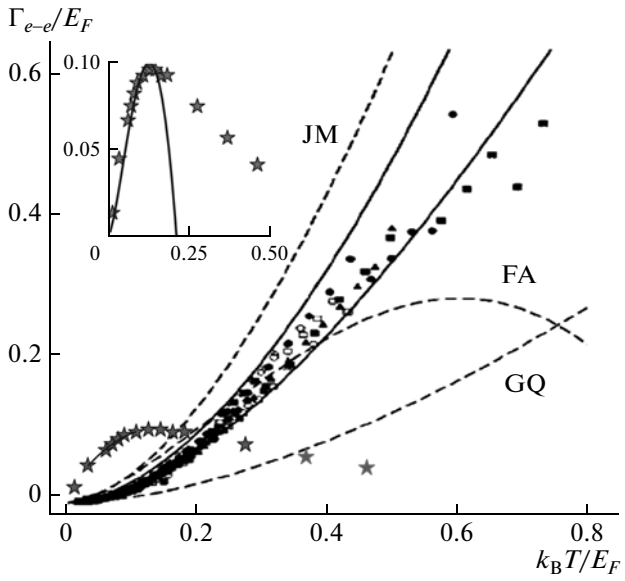

Fig. 4. This figure is taken from Turner's study [4]. The dashed lines are theoretical predictions of: **GQ** [11] ($A = 1$), **FA** [12] ($A = \pi$), and **JM** [15] ($A = \pi^2/2$). The symbols show the experimental results from [4] and the solid lines are the results of their fitting by expression (4) ($A = 3.06 \pm 0.09$) for concentrations $n = 0.3 \pm 10^{11}$ and 3 ± 10^{11} cm^{-2} (upper and lower curves, respectively). The experimental results of our study ($A = \pi^2$) are presented by asterisks. The inset shows the results of fitting our experimental data (asterisks) by theoretical expression (4) from [11] (solid line).

portional to T^n, where $n \approx 3$–5. Thus, one can suggest that the obtained temperature dependence of the inverse quantum lifetime $\sim \frac{T^2}{E_F}\ln\left(\frac{E_F}{T}\right)$ is related to only the temperature dependence of the electron–electron scattering contribution $\tau_q^{e\text{-}e}(T)$.

The dependence $\hbar/\tau_q^{e\text{-}e}(T)$ obtained by us is shown by asterisks in Fig. 4, which presents also the experimental results of [4] and a number of theoretical dependences from [11–18] (see the caption). It can be seen that the dependence $\Gamma_{e-e} = \hbar/\tau_q^{e\text{-}e}(T)$ differs radically (both qualitatively and quantitatively) from the other ones. First, it is nonmonotonic and, second, at small values of the parameter the numerical values of the electron–electron scattering rate significantly exceed both the experimental and theoretical data of other researchers.

The contribution of electron–electron scattering has been considered in many theoretical [9–20] and experimental [1–8] studies (see also references in the brief review [6]) during the last 20 years. For an ideal two-dimensional system the theoretical temperature dependence of the inverse quantum lifetime (electron–electron scattering rate) in a singlet channel in the ballistic mode was derived in [11]:

$$\frac{\hbar}{\tau_q^{e\text{-}e}}(T) = \frac{E_F}{2\pi}\left(\frac{k_BT}{E_F}\right)^2\left[\ln\left(\frac{E_F}{k_BT}\right) + \ln\left(\frac{2q_{TF}}{k_F}\right) + 1\right], \quad (4)$$

where q_{TF} is the Thomas–Fermi screening wave vector.

Turner et al. [4] processed more than 300 experimental dependences to establish that a coefficient $A = 3.06 \pm 0.09$ must be introduced into expression (4) to obtain agreement with the theory [11]. All theoretical dependences in Fig. 4 differ by only the value of factor A. To have quantitative coincidence of our experimental results with the theory at $k_BT/E_F < 0.1$, the A value must be π^2. It can be clearly seen (Fig. 4) that the discrepancy between the dependences increases with an increase in the parameter k_BT/E_F. The difference in the A values obtained by different researchers can be the result of several reasons. First, all theories were developed within perturbation theory on the assumption that $k_BT/E_F \ll 1$. Second, it was believed for a long time that two-dimensional structures, by analogy with the quantum Hall effect, are characterized by the universal dependence of the quantum lifetime on temperature [12],

$$\frac{\hbar}{\tau_q^{e\text{-}e}}(T) = A\frac{E_F}{2\pi}\left(\frac{k_BT}{E_F}\right)^2\left[\ln\left(\frac{E_F}{k_BT}\right)\right], \quad (5)$$

and it is not necessary to take into account the peculiarities of specific electron structures and experimental techniques.

However, it turned out that in the case of quasi-two-dimensional systems (in particular, single quantum wells with two quantum-confinement subbands and double quantum wells existing in equilibrium with subbands of symmetric and antisymmetric states, where electron–electron interaction occurs both within the subbands and between them), the expressions for $(\tau_q^{e\text{-}e}(T))^{-1}$ are significantly complicated [8]. Clear confirmation of this statement can be found in [18], where higher orders of expansion in Coulomb interaction and the specificity of screening in a system with two quantum wells were taken into account. Note that this complication did not lead to better agreement between the theory [11–18] and experimental data [1–8].

We revealed even more significant (not only quantitative but also qualitative) discrepancy with the existing theoretical dependences. Our experimental dependence is nonmonotonic, with an extremum observed at k_BT/E_F values of about 0.1 (see Fig. 4). A question regarding the nature of the observed feature in the temperature dependence $(\tau_q^{e\text{-}e}(T))^{-1}$ arises.

First of all, we should note that the nonmonotonic character of the dependence is present in the analyti-

cal form of function (5). All theoretical dependences $\hbar/\tau_q^{e\text{-}e}(T)$ that are known in published works [11–20] can be presented in the form

$$y = Ax^2(\ln x^{-1} + \ln B \pm C) = Ax^2\left(\ln\left[\frac{B}{x}e^{\pm C}\right]\right), \quad (6)$$

where the argument is $x = k_B T/E_F$. The coefficients A, B, and C have different analytical forms in the studies of different researchers. This circumstance is related to the choice of different models and approximations for taking into account electron–electron scattering, screening, and the peculiarities of the specific electron structure. It can easily be seen that this function also has an extremum for certain relations between the parameters. The position of the maximum depends strongly on B and C and its amplitude is determined by the parameters A, B, and C.

A values differing by an order of magnitude were reported in [11–18]: $A = \pi/4$, $1/(2\pi)$, $\pi/2$, and $\pi^2/2$ (see the caption to Fig. 4). In particular, consideration of higher expansion orders in [18] yielded $\pi/8$ and $\pi/32$ for single and double quantum wells, respectively. The situation with coefficients B and C is also ambiguous [8]. The form of the second and third terms in expression (6) is determined by the peculiarities of screening effects and specific experimental structures.

We fitted the experimental dependence (Fig. 4, inset) using expression (6). It was found that the position and amplitude of the maximum can be described by the following values of the parameters: $A = 10$, $B = 0.1$, and $C = 0$. Note that $B = 2q_{TF}/k_F = 0.1$ (according to expression (4)); this is not a physical result. For a two-dimensional electron gas in real nanostructures with a low carrier concentration, the wave vector of Thomas–Fermi screening is larger than the Fermi wave vector ($B > 1$).

Currently, the mechanism of electron–electron scattering leading to a nonmonotonic temperature dependence of the quantum lifetime is unknown. It is not inconceivable that the broadening of the peak of the resonance resistance $\rho_{xx}(B_\parallel, T)$ at $k_B T/E_F > 0.1$ is caused by not only relaxation processes.

3. CONCLUSIONS

Our study of the suppression of tunneling in a DQW in a magnetic field (parallel to the sample plane and current direction) in the temperature range $T =$ 1.8–70.0 K yielded temperature dependences of the quantum lifetime that are determined by inelastic electron–electron scattering. The obtained dependence $\tau_q^{e\text{-}e}(T)$ cannot be described within existing theories and, apparently, the observed contradiction is not due to electron momentum-relaxation processes. Possible reasons for both the qualitative and quantitative deviation of the experimental temperature dependence of electron–electron scattering from the theoretical dependence in the entire temperature range are discussed.

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